The strong coupling effect and auxiliary fields in the DGP-model

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Abstract

The DGP-model with additional terms in the action is considered. These terms have a special form and include auxiliary scalar fields without kinetic terms, which are non-minimally coupled to gravity. The use of these fields allows one to exclude the mode, which corresponds to the strong coupling effect, from the theory. Effective four-dimensional theory on the brane appears to be the same, as in the original DGP-model.

1 Introduction

The model proposed in [1] possesses a remarkable feature – a modification of gravity at ultra-large scales. This effect is very interesting from the cosmological point of view, namely for possible explanation of the observed large-scale acceleration of our Universe [2]. Unfortunately the strong coupling effect due to the existence of a strongly interacting mode, which is a four-dimensional scalar, was found in the model [3, 4]. This problem was discussed in a lot of papers and there were proposed some mechanisms of its solution. For example, it was argued that the strong coupling effect is an artifact of linear approximation and disappears after taking into account the non-linear effects (resummation of perturbative series) [5]. Another approach is to utilize a regularization in the model [6, 7]. In this paper a new mechanism for overcoming the strong coupling, based on a slight modification of the original action, is proposed.

The paper is organized as follows. First, we solve exactly equations for linearized gravity in the DGP-model and explicitly distinguish the strong coupled mode, which is nothing else but the h_{44}-component of metric fluctuations. This result is well-known and was mentioned earlier, for example, in [8]. Then we show that with the help of additional terms in the action, which contain "auxiliary" scalar fields non-minimally interacting with gravity, one can exclude this strongly interacting mode from the theory. This is shown by solving the corresponding equations of motion and identifying the physical degrees of freedom of the model. At the same time effective equation for the $\mu\nu$-component of metric fluctuations remains unchanged and one can use all phenomenological predictions obtained for gravity on the brane. Finally we briefly discuss the obtained results and possible consequences.

2 Gravity in the original DGP-model

Action of the model proposed in [1] has the following form:

$$S_{DGP} = M_*^3 \int R\sqrt{-g}\, d^4x dy + \Omega M_*^3 \int_{y=0} \tilde{R}\sqrt{-\tilde{g}}\, d^4x, \qquad (1)$$

where M_* is the five-dimensional Planck mass, $\Omega >> 1/M_*$ and $\tilde{g}_{\mu\nu}$ is the induced metric on the brane, which is located at the point $y = 0$ of the extra dimension. In addition, let us suppose that the model possesses $y \leftrightarrow -y$ symmetry, which fixes the brane position. We also note that the signature of the metric g_{MN} is chosen to be $(-, +, +, +, +)$.

Let us denote $\hat{\kappa} = \frac{1}{\sqrt{M_*^3}}$ and parameterize the metric g_{MN} as

$$g_{MN} = \gamma_{MN} + \hat{\kappa} h_{MN}, \tag{2}$$

h_{MN} being the metric fluctuations and $M, N = 0, .., 4$. Substituting this parameterization into (1) and retaining the terms of the zero order in $\hat{\kappa}$, we can get the second variation action of the model. This action is invariant under the gauge transformations

$$\begin{aligned} h'_{\mu\nu}(x,y) &= h_{\mu\nu}(x,y) - (\partial_\mu \xi_\nu + \partial_\nu \xi_\mu), & (3)\\ h'_{\mu 4}(x,y) &= h_{\mu 4}(x,y) - (\partial_\mu \xi_4 + \partial_4 \xi_\mu), & (4)\\ h'_{44}(x,y) &= h_{44}(x,y) - 2\partial_4 \xi_4, & (5) \end{aligned}$$

and the functions $\xi^M(x,y)$ satisfy the symmetry conditions

$$\begin{aligned} \xi^\mu(x,-y) &= \xi^\mu(x,y), \\ \xi^4(x,-y) &= -\xi^4(x,y). \end{aligned} \tag{6}$$

Here $\mu, \nu = 0, .., 3$. The matter is assumed to be placed on the brane only, and its interaction with gravity has the standard form

$$\frac{\hat{\kappa}}{2} \int_{y=0} h^{\mu\nu}(x,0) t_{\mu\nu}(x) d^4x, \tag{7}$$

$t_{\mu\nu}(x)$ denoting the energy-momentum tensor of the matter.

It is necessary to emphasize that in general all fluctuations of metric must satisfy the physical boundary conditions at $y \to \pm\infty$, $x^i \to \pm\infty$ $(i = 1, 2, 3)$, i.e. vanish at spatial infinity. This is a reasonable assumption - for example, the h_{00}-component is associated with Newton's potential, which must vanish at infinity (for the matter, which is localized in some finite domain). Obviously, the gauge functions $\xi^M(x,y)$ must be finite everywhere – it follows simply from the definition of the gauge transformations. It means that $\xi_M(x,y)$ must be finite too (the problem of physical boundary conditions in the case of RS2 model was discussed in detail in [9]).

The equations of motion for the metric fluctuations, corresponding to action (1), have the following form (they can be easily derived from the corresponding equations in [9, 10]):
1) $\mu\nu$-component

$$\begin{aligned} &\frac{1}{2}\left(\Box h_{\mu\nu} - \partial_\mu \partial^\rho h_{\rho\nu} - \partial_\nu \partial^\rho h_{\rho\mu} + \partial_4 \partial_4 h_{\mu\nu} + \partial_\mu \partial_\nu h + \partial_\mu \partial_\nu h_{44}\right) - \\ -\ &\frac{1}{2}\partial_4(\partial_\mu h_{\nu 4} + \partial_\nu h_{\mu 4}) + \frac{1}{2}\gamma_{\mu\nu}\left(\partial^\rho \partial^\sigma h_{\rho\sigma} - \Box h - \partial_4 \partial_4 h - \Box h_{44} + 2\partial^\rho \partial_4 h_{\rho 4}\right) + \\ +\ &\frac{\Omega}{2}\delta(y)\left[\left(\Box h_{\mu\nu} - \partial_\mu \partial^\rho h_{\rho\nu} - \partial_\nu \partial^\rho h_{\rho\mu} + \partial_\mu \partial_\nu h\right) + \gamma_{\mu\nu}\left(\partial^\rho \partial^\sigma h_{\rho\sigma} - \Box h\right)\right] = \\ =\ &-\frac{\hat{\kappa}}{2} t_{\mu\nu}(x)\delta(y), \end{aligned} \tag{8}$$

2) $\mu 4$-component

$$\partial_4(\partial_\mu h - \partial^\nu h_{\mu\nu}) - \partial^\nu \left(\partial_\mu h_{\nu 4} - \partial_\nu h_{\mu 4}\right) = 0, \tag{9}$$

3) 44-component

$$\partial^\mu \partial^\nu h_{\mu\nu} - \Box h = 0, \tag{10}$$

where $h = \eta^{\mu\nu} h_{\mu\nu}$ and $\Box = \eta^{\mu\nu}\partial_\mu \partial_\nu$.

In what follows, we will also use an auxiliary equation, which is obtained by substituting the equation for 44-component into the contracted equation for $\mu\nu$-component. This equation contains h, $h_{\mu 4}$ and h_{44} only and has the form:

$$\partial_4 \partial_4 h - 2\partial_4 \partial^\mu h_{\mu 4} + \Box h_{44} = \frac{\hat{\kappa}}{3}\, t^\mu_\mu(x)\delta(y), \tag{11}$$

where $\Box = \eta^{\mu\nu}\partial_\mu \partial_\nu$. By integrating this equation over y in the limits $(-\infty, \infty)$ and using the physical boundary conditions for the fields $h_{\mu\nu}$, $h_{\mu 4}$ and h_{44}, we find that the function $\rho(x)$, defined by

$$\rho(x) = \int_{-\infty}^{\infty} h_{44}(x, y)\, dy, \tag{12}$$

is not equal to zero and satisfies the equation

$$\Box \rho(x) = \frac{\hat{\kappa}}{3}\, \eta^{\mu\nu} t_{\mu\nu}(x) \equiv \frac{\hat{\kappa}}{3}\, t(x). \tag{13}$$

This means that the field h_{44} cannot be gauged out, because otherwise the equations of motion for linearized gravity become inconsistent. One can easily check that in the absence of the fields $h_{\mu 4}$ and h_{44} from equation (11) follows that

$$h \sim t|y|. \tag{14}$$

Thus h diverges at $y = \pm\infty$ and does not satisfy the physical boundary conditions.

We will use the following form of ξ_4 to impose an appropriate gauge on the field h_{44} [9]:

$$\xi_4(x, y) = \frac{1}{4}\int_{-y}^{y} h_{44}(x, y')dy' - \frac{1}{4C}\int_{-y}^{y} F(y')dy' \int_{-\infty}^{\infty} h_{44}(x, y')dy', \tag{15}$$

where $F|_{y\to\pm\infty} = 0$ and

$$C = \int_{-\infty}^{\infty} F(y)dy. \tag{16}$$

Note that ξ_4 satisfies the symmetry and the boundary conditions. With the help of (15) we can pass to the gauge, in which

$$h_{44}(x, y) = F(y)\phi(x), \tag{17}$$

where

$$\phi(x) = \frac{1}{C}\int_{-\infty}^{\infty} h_{44}(x, y)dy \tag{18}$$

and depends on x only (it is evident that $\rho(x) = C\phi(x)$). It turns out to be convenient to choose $F(y) = e^{-k|y|}$. Obviously, the field h_{44} satisfies the symmetry and the physical boundary conditions in this gauge. Moreover, we have no residual gauge transformations with ξ_4. We also note that since $\xi_4(x,0) = 0$, the brane remains unshifted in this gauge.

Now let us discuss the gauge condition for the field $h_{\mu 4}$. Let us take the gauge function $\xi_\mu(x,y)$ in the following form:

$$\xi_\mu(x,y) = \int_{-\infty}^{y} h_{\mu 4}\left(x, y'\right) \, dy'. \tag{19}$$

One can easily see that due to the symmetry $h_{\mu 4}(x,-y) = -h_{\mu 4}(x,y)$, $\xi_\mu(x,y)$ satisfies the symmetry condition $\xi_\mu(x,-y) = \xi_\mu(x,y)$. Moreover, it is easy to see that $\xi_\mu(x,y)|_{y\to\pm\infty} \to 0$, at least in the sense of the principal value of the integral in equation (19) (again due to the symmetry of $h_{\mu 4}$), i.e. it is finite. Finally, it is not difficult to check that the gauge transformation with ξ_μ given by (19) gauges the field $h_{\mu 4}$ out. It seems that this formal argumentation can be used in favor of the possibility to make the $h_{\mu 4}$-field vanish everywhere. Anyway, as we will see later, equations of motion can be solved in the gauge $h_{\mu 4} = 0$

After this gauge fixing we are still left with residual gauge transformations of the form

$$\partial_4 \xi_\mu = 0. \tag{20}$$

Using the physical boundary conditions for the field $h_{\mu\nu}$ at $y \to \infty$, from equation (3) it follows that

$$\xi_\mu = 0. \tag{21}$$

Thus, we have no residual gauge transformations with $\xi_\mu \neq 0$.

The substitution, which allows us to solve equations of motion in the gauge $h_{\mu 4}(x,y) = 0$, $h_{44}(x,y) = e^{-k|y|}\phi(x)$, has the form

$$h_{\mu\nu} = b_{\mu\nu} - \frac{1}{k^2} e^{-k|y|} \partial_\mu \partial_\nu \phi. \tag{22}$$

Substituting (22) into (9), (10), (11) we get

$$\partial_4(\partial_\mu b - \partial^\nu b_{\mu\nu}) = 0, \tag{23}$$

$$\partial^\mu \partial^\nu b_{\mu\nu} - \Box b = 0, \tag{24}$$

$$\partial_4 \partial_4 b + \frac{2}{k} \Box \, \phi \delta(y) = \frac{\hat{\kappa}}{3} t \delta(y). \tag{25}$$

Integrating (25) over y in the limits $(-\infty, \infty)$ and using the physical boundary conditions for the field $b_{\mu\nu}$, we get

$$\Box \, \phi = \frac{\hat{\kappa} k}{6} t. \tag{26}$$

This mean that

$$\partial_4 b = B(x), \tag{27}$$

where $B(x)$ is some function of x only. Using the symmetry conditions, we obtain $B(x) \equiv 0$.

From the physical boundary conditions it follows that

$$b = 0, \tag{28}$$

$$\partial^\mu b_{\mu\nu} = 0. \tag{29}$$

where $b = \eta^{\mu\nu} b_{\mu\nu}$. It is not difficult to find corresponding equation for the field $b_{\mu\nu}$.

It is evident that such coupling constant for the field ϕ makes $\hat{\kappa} h_{44}$ to be much larger than unity if $\Omega >> 1/M_*$ (for example, in the case $M_* \sim 10^{-3} eV$) at least at $y = 0$. Linearized approximation breaks down, and this is the origin of the so-called strong coupling effect.

3 Modified DGP-model

To solve this problem let us add to the action (1) the following terms

$$S_{add} = M_*^3 \int_{y=0} \varphi(x) \tilde{R} \sqrt{-\tilde{g}}\, d^4x + \int \Phi(x,y) R^2 \sqrt{-g}\, d^4x dy. \tag{30}$$

The first term in (30) corresponds to the four-dimensional Brans-Dicke theory with $\omega = 0$ (where ω is the Brans-Dicke parameter). The absence of the kinetic terms for the fields $\varphi(x)$ and (especially) $\Phi(x,y)$ look rather strange. Nevertheless one can recall that SUSY is based on the use of such "auxiliary" fields, which are necessary for implementing the supersymmetry transformations. A simple example with the fields of such type in classical field theory can be found in [11]. In any case, there are no strong objections against considering fields of this type, although these fields will be used for other purposes, not that as in SUSY. We also assume that $\varphi(x) \equiv 0$ in the background, otherwise the redefinition of Ω in equation (1) is needed.

From the equation of motion for the field $\Phi(x,y)$ we have the following condition for the five-dimensional curvature

$$R = 0. \tag{31}$$

This equation holds in the bulk.

Equations (8), (9), (10) and (11) are modified by the terms of (30) and take the form (it is easy to see that one can use the same gauge condition for the field $h_{\mu 4}$ as that used in the previous section):

1) $\mu\nu$-component

$$\begin{aligned} &\frac{1}{2}\left(\Box h_{\mu\nu} - \partial_\mu \partial^\rho h_{\rho\nu} - \partial_\nu \partial^\rho h_{\rho\mu} + \partial_4 \partial_4 h_{\mu\nu} + \partial_\mu \partial_\nu h + \partial_\mu \partial_\nu h_{44}\right) + \\ +\; &\frac{1}{2}\gamma_{\mu\nu}\left(\partial^\rho \partial^\sigma h_{\rho\sigma} - \Box h - \partial_4 \partial_4 h - \Box h_{44}\right) + \\ +\; &\frac{\Omega}{2}\delta(y)\left[\left(\Box h_{\mu\nu} - \partial_\mu \partial^\rho h_{\rho\nu} - \partial_\nu \partial^\rho h_{\rho\mu} + \partial_\mu \partial_\nu h\right) + \gamma_{\mu\nu}\left(\partial^\rho \partial^\sigma h_{\rho\sigma} - \Box h\right)\right] + \\ +\; &\frac{1}{\hat{\kappa}}\left(\partial_\mu \partial_\nu \varphi - \eta_{\mu\nu} \Box \varphi\right)\delta(y) = -\frac{\hat{\kappa}}{2} t_{\mu\nu}(x)\delta(y), \end{aligned} \tag{32}$$

(corresponding equations for the Brans-Dicke theory can be found, for example, in [12]),

2) $\mu 4$-component

$$\partial_4(\partial_\mu h - \partial^\nu h_{\mu\nu}) = 0, \tag{33}$$

3) 44-component

$$\partial^\mu \partial^\nu h_{\mu\nu} - \Box h = 0, \tag{34}$$

4) auxiliary equation

$$\partial_4 \partial_4 h + \Box h_{44} + \frac{2}{\hat{\kappa}} \Box \varphi \delta(y) = \frac{\hat{\kappa}}{3} t(x) \delta(y), \tag{35}$$

The second term in (30) does not contribute to the equations since it contains R squared. Thus the main purpose of this term is to make the five-dimensional curvature equal to zero in the bulk. It is evident that equation of motion for the field $\varphi(x)$ does not contradict (34). Thus only the case $\omega = 0$ is compatible with action (1), otherwise the existence of kinetic term for $\varphi(x)$ leads to the condition $\Box \varphi = 0$ and this field decouples from the theory. Linearizing equation (31) and taking into account (34), we get

$$\partial_4 \partial_4 h + \Box h_{44} = 0. \tag{36}$$

It means that

$$\Box \varphi = \frac{\hat{\kappa}^2}{6} t. \tag{37}$$

Integrating (36) in the gauge $h_{44}(x, y) = e^{-k|y|} \phi(x)$ in the limits $(-\infty, \infty)$ and using the physical boundary conditions for the field $h_{\mu\nu}$, we get

$$\Box \phi = 0. \tag{38}$$

We see, that h_{44}-field does not interact with matter on the brane, i.e. it decouples from the effective theory. This also means that the linear approximation is valid. Moreover, we can totally gauge out this field with the help of ξ_4 of the form

$$\xi_4(x, y) = \frac{1}{4} \int_{-y}^{y} h_{44}(x, y') dy' \tag{39}$$

instead of (15). Analogously to what was made in the previous section, we have the following conditions for the field $h_{\mu\nu}$:

$$h = 0, \tag{40}$$

$$\partial^\mu h_{\mu\nu} = 0. \tag{41}$$

Let us explain what has happened. The field φ takes over the role of the h_{44}-component of metric fluctuations. Since this field does not exist "inside" the curvature, one does not need to worry about its absolute value – it does not affect the validity of linear approximation. At the same time the field Φ makes h_{44} not interacting with matter on the brane (one can check that in the absence of the corresponding term it is not necessarily so).

Now we are ready to find the equation for the $\mu\nu$-component. In gauge (40), (41) and subject to (37) it takes the well-known form

$$\Box h_{\mu\nu} + \partial_4 \partial_4 h_{\mu\nu} + \Omega \delta(y) \Box h_{\mu\nu} = -\hat{\kappa} \delta(y) \left(t_{\mu\nu} - \frac{1}{3} \left[\eta_{\mu\nu} - \frac{\partial_\mu \partial_\nu}{\Box} \right] t \right). \tag{42}$$

It coincides with that derived in [1] and was examined in a lot of papers, so we will not discuss it here. All the predictions obtained with the help of this equation are valid in this case too.

4 Discussion and final remarks

In this paper we discussed the DGP-model with additional terms, which include two "auxiliary" scalar fields. The first one takes over the role of the h_{44}-component of metric fluctuations (the radion field), which is indispensable for the original DGP-model and makes the equations consistent, whereas the second one is responsible for dropping out this component from the theory. Thus, the strongly interacting scalar mode vanishes away as well as the strong coupling effect. The effective equation for the $\mu\nu$-component of metric fluctuations appears to be the same, as in the case of the original DGP setup. Also it is not difficult to check that non-linear term $\sim \varphi \tilde{R}_{\mu\nu}$, which appears in the Einstein equations, can be neglected in comparison with the term $\sim \Omega \tilde{R}_{\mu\nu}$ even for such massive objects like those in our Solar system.

Now let us discuss the problem of ghosts in the model. First, we note, that the four-dimensional Brans-Dicke theory with $\omega = 0$ is stable, which can be shown with the help of conformal rescaling. It is reasonable to suppose, that this property survives in the five-dimensional theory too. Second, the field $\varphi(x)$ behaves exactly as the radion at the linear order (compare (23), (24), (25) and (33), (34), (35)), which is not a ghost in the original DGP-model. And third, the existence of ghosts lead to additional repulsive forces, which are absent in the case under consideration. This formal argumentation can be used in favor of the absence of ghosts at least in the second variation Lagrangian.

In this paper we solve equations of motion only at the linear order. One can ask about the possibility of the presence of the strong coupling effect at the next orders, for example, through the cubic and quartic terms of the effective Lagrangian. To this end it is necessary to make a thorough analysis analogous to the one, which was made in [4]. Nevertheless the strong coupling effect manifests itself even at the linear order, as it was shown in Section 2. This can be also seen through the diagonalization of the second variation Lagrangian of the model. In the original DGP-model such procedure leads to appearance of the kinetic term for the radion with a small normalization constant, leading to the strong coupling at the next orders (see [3, 4]). As it was shown above, the field $\varphi(x)$ acts in the same way as the radion (at the linear order). Nevertheless, these fields have different nature, which becomes apparent after diagonalization of the Lagrangian. In the four-dimensional theory diagonalization of the Brans-Dicke action by conformal rescaling results in the appearance of the scalar field, interacting directly with matter with the same coupling constant as the tensor gravity. The linearized substitution for the $\mu\nu$-component of metric fluctuations, which corresponds to the conformal rescaling, does not contain large pieces, which can lead to the strong coupling at the next orders. In the case of the model described in Section 3 the corresponding diagonalization is not so trivial, as that in the four-dimensional case. Nevertheless naive calculations show that situation appears to be analogous to the four-dimensional case: the main contributions to the kinetic term of the field $\varphi(x)$ come from the brane action, and these contributions appear to be much larger than those coming from the five-dimensional curvature. At the same time the field $\varphi(x)$ can appear, for example, in the cubic terms of the action only through the substitution for the $\mu\nu$-component, which seems not to contain large pieces, contrary to the case of original DGP-setup, in which analogous substitution does contain large pieces (see [4]). Moreover, the radion also "lives" in the five-dimensional curvature and can appear in the cubic order terms "independently"

of the substitution. In other words, the four-dimensional Brans-Dicke gravity with $\omega = 0$ dominates on the brane at the small distances, instead of the four-dimensional gravity with contribution of the radion in the original DGP-model. If there were no radion, which leads to the strong coupling, in the original DGP-setup, the only pure four-dimensional gravity would live on the brane at the small distances. After the modification of the model the radion drops out from the theory, whereas the four-dimensional Brans-Dicke gravity takes over the role of the "pure four-dimensional gravity" of the original DGP-model. Thus, there seem to be no prerequisites for the strong coupling effect, coming from the second variation Lagrangian, contrary to the case of the original DGP-model.

However, since the absence of the strong coupling effect was shown only at the level of solving the classical linear equations of motion, a further thorough analysis is necessary. This analysis can also clarify the fate of the radion and the $\mu 4$-component of metric fluctuations in the model.

One can worry about the fact that the field Φ is not defined by the equations of motion. This problem can be easily solved by adding an extra term to the action, for example, one can choose

$$S_\Phi = \frac{\alpha}{3} \int \Phi^3(x,y)\sqrt{-g}\, d^4x dy, \tag{43}$$

where α is real and positive. In this case equation (31) will take the following form:

$$R^2 + \alpha\Phi^2 = 0, \tag{44}$$

which uniquely determines the field Φ ($\Phi \equiv 0$).

Here we do not consider the problem of the incorrect tensor structure of the graviton, which is analogous to the vDVZ-discontinuity in massive gravity [13, 14]. This issue was discussed in a lot of papers including [1]. It seems that this problem is inherent at least to five-dimensional models with flat background and matter localized on the brane (see equation (34)). It is reasonable to suppose that it can be solved in models with warped geometry. Hence the modification of the DGP-model proposed above does not make the model applicable for describing gravity in our world. Nevertheless one can hope that the use of additional fields non-minimally coupled to gravity (not necessarily without kinetic terms) can be quite efficient in some more realistic models, which admit long-range modification of gravity and suffer from the problem of strong coupling effect.

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